Chiral Extrapolation, Renormalization, and the Viability of the Quark Model

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The relationship of the quark model to the known chiral properties of QCD is a longstanding problem in the interpretation of low energy QCD. In particular, how can the pion be viewed as both a collective Goldstone boson quasiparticle and as a valence quark antiquark bound state where universal hyperfine interactions govern spin splittings in the same way as in the heavy quark systems. We address this issue in a simplified model which; however, reproduces all features of QCD relevant to this problem. A comparison of the many-body solution to our model and the constituent quark model demonstrates that the quark model is sufficiently flexible to describe meson hyperfine splitting provided proper renormalization conditions and correct degrees of freedom are employed consistently.

One of the fundamental symmetries of the strong interaction is chiral symmetry (χS). This symmetry arises because the bare masses of the light u and d quarks are small with respect to hadronic scales (or, $m_{u,d} << \Lambda_{QCD}$). The nonzero quark masses imply that chiral symmetry is explicitly broken. Chiral symmetry is also broken spontaneously by the QCD vacuum. This fact is represented by the nonzero value of the quark condensate which plays the role of an order parameter.

The phenomena associated with chiral symmetry breaking, for example the spectral properties of quarks in the dressed vacuum or the formation of Goldstone modes, are among the most important and most extensively studied aspects of strong QCD. As a consequence of chiral symmetry, the axial current is (partially) conserved and the interactions of low momentum pions with hadrons are weak. This forms the basis of chiral perturbation theory and, more generally, enables one to formulate an effective approach to low energy QCD with dynamics being dominated by weakly interacting Goldstone bosons and nucleons [1]. At high temperatures or densities the chiral invariance of the QCD vacuum is expected to be restored. The thermodynamics of this phenomenon plays an important role in determining signatures of the quark–gluon plasma in relativistic heavy ion collisions [2]. Finally, χS is also relevant for lattice gauge simulations. In particular, chiral symmetry predicts a highly nonlinear dependence of a number of observables on the light quark mass. Since current lattice QCD simulations are limited to light quarks with masses of the order of 200 MeV the interpretation of lattice results requires input from chiral symmetry-based phenomenology. It is clear that chiral symmetry plays a central role in determining the structure and interactions of low energy hadrons [3].

The constituent quark model (CQM) has historically developed in parallel with chiral theory and has been extensively used as a simple alternative for the study of hadronic phenomenology. Even though it is loosely connected to QCD, the quark model has had a large number of remarkable successes [4]. The CQM is based on the idea that strong interactions lead to massive quasiparticles (constituent quarks) and that hadronic structure is dominated by the interactions between valence constituent quarks. One argues that constituent quarks are the effective degrees of freedom arising after dynamical χS breaking due to bare quark interactions with the quark sea in the chiral noninvariant vacuum [5]. In the CQM, however, properties of the constituent quarks, *e.g.* masses and magnetic moments, are treated as free parameters making the approach effectively insensitive to the underlying chiral structure of QCD. This implies, in particular, that CQM pions lose their nature as Goldstone bosons and are not much different from, say, ρ mesons. The mass splitting between the pion (a spin-0 constituent $Q\bar{Q}$ state) and the ρ meson (a spin-1 state) is attributed to a residual hyperfine interaction. The hyperfine interaction is typically associated with the nonrelativistic reduction of single gluon exchange between constituent quarks and yields a vector–pseudoscalar splitting proportional to $\alpha_s/(m_q m_{\bar{q}})$.

Since the notion of gluon exchange is intrinsically perturbative, it is likely a useful concept in heavy quark systems where the large quark mass guarantees its applicability. However its relevance to the light quark sector is unclear. Indeed, one may expect that the π-ρ mass splitting is largely driven by the underlying χS rather than by the hyperfine interaction. It is therefore surprising that when the splitting between the flavored quark-antiquark vector and pseudoscalar mesons is plotted against the mass of one of the constituent quarks (the other constituent quark mass is held fixed), it does indeed behave like $1/m_{con}$ typical to a hyperfine interaction all the way down to the light quark sector [6]. This is shown by the squares in Fig. 1. The definition of the CQM masses depends on details of the model, here we use typical values $m_u = m_d = 330$ MeV, $m_s = 550$ MeV, $m_c = 1600$ MeV, and $m_b = 4980$ MeV.

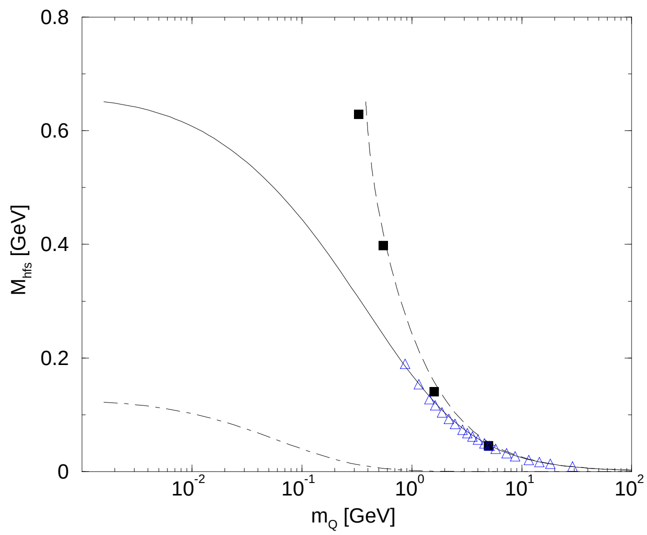

FIG. 1. Flavored pseudoscalar and vector meson mass splittings as a function of the renormalized heavy quark mass (solid, $m_Q = m_R$) and constituent quark mass (dashed, $m_Q = m_{con}$) together with lattice results [12] (triangles, $m_Q = m_R$) and experiment [13] (squares, $m_Q = m_{con}$). The dot-dash curve represents the hyperfine splitting obtained when the hyperfine interaction is neglected. The renormalized light quark mass is fixed at 5 MeV.

The question we address here is whether the apparent universality of the effective interactions of the constituent quark model is consistent with the dictates of chiral symmetry in the light meson sector. To address this issue it is natural to work within the framework of Hamiltonian Coulomb gauge QCD. This enables us to map the degrees of freedom and the global symmetries of QCD to those of the constituent quark model and thus effectively "derive" the quark model from QCD. The Coulomb gauge QCD Hamiltonian can be formally written as [7,8] $H = H_{can}(\Lambda) + \delta H(\Lambda)$. Here the first term is the canonical Hamiltonian which contains the Coulomb interaction $H_{Coulomb}$ regularized at a scale Λ. The cutoff Λ is necessary since products of field operators at the same point lead to divergences. Modifying the structure of the canonical Hamiltonian by the cutoff requires counterterms, δH which are to be adjusted so that physical observables are Λ-independent. We note that in the chiral limit Λ is the only scale in the Hamiltonian and that light meson masses will therefore be proportional to the cutoff.

Even though one can in principle choose an arbitrary value for the cutoff, there is typically an optimal choice, $\Lambda = \Lambda_R$ which depends on the basis and approximation scheme used to diagonalize the Hamiltonian. Loosely speaking, as Λ increases beyond Λ_R, the strength of the various interaction terms decrease but at the same time the available phase space for partons to interact increases. On the other hand, for small cutoffs the phase space decreases but the remaining Fock sectors interact strongly, making it difficult to determine the effective degrees of freedom. Since our goal is to understand the quark model from QCD, we choose Λ_R to match to the quark model. The quark model scale is the scale where all QCD dynamics are given in terms of effective interactions between valence constituent quarks (that is if the quark model is to be a rigorous consequence of QCD). Thus, in particular, since soft gluons are not present in the Fock space at the quark model scale there are no contributions to hyperfine interactions from soft gluon intermediate states.

We now examine the appropriate Λ_R for heavy-heavy and heavy-light systems. Since the heavy-heavy case is analogous to Coulombic bound states, it is natural to choose Λ_R so that it scales with $m_R \equiv m(\Lambda = \Lambda_R)$ – the heavy quark mass renormalized at the scale Λ_R. This constrains the average momenta of heavy quarks to lie below m_R. Furthermore, since interactions between heavy quarks and transverse gluons vanish in the nonrelativistic limit, gluons and heavy quarks decouple and the dominant interaction between heavy quarks is due to the nonabelian instantaneous Coulomb potential. As discussed in detail in Ref. [9], in nonrelativistic QCD (NRQCD) the proper scaling of the quark-gluon interaction is obtained if $\Lambda_R = \alpha_R m_R$ where $\alpha_R = \alpha(\Lambda_R) = \langle p \rangle / m_R$. Here $\langle p \rangle$ denotes the typical momentum of a heavy quark in the bound state. The scaling $\langle p \rangle \sim \alpha_R m_R$ and $E - 2m_R \sim \alpha_R^2 m_R$ for average momenta and energies is typical for a Coulombic bound state.

We now consider the heavy-light case. In this case the dynamics of the "brown muck" of gluons and light quarks is still subject to the full complexity of nonperturbative strong interactions. A heavy-light quark system with a single heavy quark is no longer Coulombic and thus the average separation between the quarks increases causing them to

probe the long range confinement potential. Quark model phenomenology [10] dictates that this potential be linear. Thus the relevant quantities are the scales of the linear potential b and the kinetic energy, which is characterized by the reduced mass μ. Dimensional analysis then indicates that the average momenta in the bound state scale as $\langle p \rangle = (b\mu)^{1/3}$. Since the reduced mass is very nearly equal to the light constituent quark mass, and this mass and $\sqrt{b}$ are of the order of Λ_{QCD}, we simply use $\Lambda_R = \langle p \rangle \approx \Lambda_{QCD}$ for the heavy-light case.

Proper determination of the renormalization scale is the first step toward "derivation" of the quark model from QCD. The second, is to implement dynamical chiral symmetry breaking which will generate masses for the light bare quarks thereby creating the constituent quarks of the CQM. The constituent Fock space then provides an efficient basis for diagonalizing the Hamiltonian.

In this letter we do not aim at a rigorous derivation of the meson spectrum. Our goal is to illustrate how the two steps discussed above enable one to derive the constituent quark picture while maintaining correct chiral properties. We will therefore approximate H and use a simple contact interaction instead of the full nonabelian Coulomb kernel. This is sufficient as long as the interaction causes the typical momenta in the bound state to scale as $\langle p \rangle = \alpha_R m_R$ for heavy-heavy and $\langle p \rangle = \Lambda_{QCD}$ for heavy-light mesons.

$$H_{Coulomb} \rightarrow H_c(\Lambda) = \frac{c_c(\Lambda)}{\Lambda^2} \int d\mathbf{x} \left[\psi^\dagger(\mathbf{x}) \mathbf{T} \psi(\mathbf{x}) \psi^\dagger(\mathbf{x}) \mathbf{T} \psi(\mathbf{x}) \right]_\Lambda \tag{1}$$

Here the subscript Λ indicates that the operators are to be point-split or smeared over a distance $\sim 1/\Lambda$ [8] and $c_c(\Lambda)$ is the coupling which replaces $\alpha(\Lambda)$. The counterterms δH contain relevant, marginal, and irrelevant operators. In the chiral limit, chiral symmetry prevents the occurrence of a relevant operator in the quark sector. For finite quark masses, the relevant operator is absorbed into the definition of the bare quark mass. The effect of transverse gluons eliminated by the cutoff show up through contact operators of dimension six or greater. There are a number of such terms; to illustrate the effect of hyperfine interactions we consider the dominant spin-dependent term,

$$H_h(\Lambda) = \frac{c_h(\Lambda)}{\Lambda^2} \int d\mathbf{x} \left[\psi^\dagger(\mathbf{x}) \mathbf{T} \boldsymbol{\alpha} \psi(\mathbf{x}) \psi^\dagger(\mathbf{x}) \mathbf{T} \boldsymbol{\alpha} \psi(\mathbf{x}) \right]_\Lambda , \tag{2}$$

where $c_h(\Lambda)$ is proportional $\alpha(\Lambda)$. Thus the full model Hamiltonian is given by $H = \int \psi^\dagger[-i\alpha \cdot \nabla + \beta m(\Lambda)]\psi + H_c(\Lambda) + H_h(\Lambda)$.

As stated above, generating spontaneous chiral symmetry breaking is crucial to examining the interplay of the hyperfine interaction and the Goldstone boson nature of the pion. As a result, we employ the BCS ansatz to construct an approximation to the chirally noninvariant vacuum. For the above Hamiltonian this leads to the following mass gap equation which determines the effective constituent quark mass

$$m_{con} = m(\Lambda) + \frac{m_{con}}{\Lambda^2} \left(\tilde{c}_c(\Lambda) - \tilde{c}_h(\Lambda) \right) \int^\Lambda q^2 dq \frac{1}{\sqrt{m_{con}^2 + \mathbf{q}^2}}, \tag{3}$$

where $\tilde{c}_c = C_F c_c / 2\pi^2$, $\tilde{c}_h = 3 C_F c_h / 2\pi^2$, and the single quasiparticle energies are given by $E(\mathbf{q}) = \sqrt{m_{con}^2 + \mathbf{q}^2}$. These resemble energies of the constituent quarks if $m_{con} \sim 300$ MeV for the light quarks.

We wish to carefully distinguish the three types of quark masses which have appeared. They are the bare cutoff-dependent quark mass, $m(\Lambda)$ which is defined as the parameter appearing in the mass term of the canonical Hamiltonian, the renormalized quark mass, m_R, given by $m(\Lambda_R)$, and the constituent quark mass, m_{con}, which is the nonperturbative mass obtained by solving the gap equation. The cutoff dependence of the gap equation and m_{con} are discussed later.

The RPA or Bethe-Salpeter equation for the meson bound state, M, is then given by [11]

$$\langle M | [H, Q_M^\dagger] | BCS \rangle = (E_M - E_{BCS}) \langle M | Q^\dagger | BCS \rangle, \tag{4}$$

where $Q_M^\dagger$ is defined in a standard way in terms of the positive and negative energy wave functions, $Q_M^\dagger = \sum_{\alpha\beta} \left[\psi^+_{\alpha\beta} B^\dagger_\alpha D^\dagger_\beta - \psi^-_{\alpha\beta} D_\beta B_\alpha \right]$ with B and D being the quasiparticle operators. In the simple approximation to H used here, Eq. (4) is an algebraic equation for the bound state masses $E_M = E_M(\Lambda, m(\Lambda), \tilde{c}_c(\Lambda), \tilde{c}_h(\Lambda))$. Requiring the E_M to be Λ-independent is used to determine the cutoff dependence of the couplings.

In the following we will concentrate on the pseudoscalar (ps) and vector (v) open flavor mesons. As discussed above, in the case of unequal quark masses, the renormalization constant should be fixed at $\Lambda_R \approx \Lambda_{QCD}$. Here we use the numerical value 420 MeV because this corresponds to $\sqrt{b}$. The average, renormalized light quark mass, $m_R \equiv m(\Lambda_R)$ is set to 5 MeV, and the two renormalized couplings $\tilde{c}_{R,c,h} \equiv \tilde{c}_{c,h}(\Lambda_R)$ are determined by fitting the π and ρ meson masses, $E_{\pi,\rho} = E_{\pi,\rho}(\Lambda_R, m_R, \tilde{c}_{R,c}, \tilde{c}_{R,h})$.

Solving the gap equation yields a light constituent quark mass of 380 MeV. For a heavy-light meson, the solid line in Fig. 1 shows the dependence of the hyperfine mass splitting $M_{hfs} \equiv E_v - E_{ps}$ as a function of the renormalized heavy quark mass $m_Q = m_R$ with the renormalized light quark mass fixed at 5 MeV. As m_Q increases the splitting falls off as $1/m_Q$, however, as the heavy quark mass approaches the light quark limit the slope changes, reflecting the emergence of chiral symmetry: $E_{ps} \propto m_Q^{1/2}$. Our predictions reproduce the lattice results (triangles in Fig. 1) [12] available for large quark masses. The lattice results are given for $M_{hfs} \cdot a$ as a function of $m_R \cdot a$ where a is the lattice spacing. Lattice calculations match our results if one takes $a^{-1} = 1.45$ GeV which is very close to that of Ref. [14]. It is nontrivial that our predictions match the lattice over a large range of quark masses since the coupling constants were fixed in the chiral limit. We have fit the ansatz $M_{hfs} = a + b\sqrt{m_R + d} + cm_R$ to the data of Fig. 1. This ansatz has been employed in Ref. [15] to fit their lattice data and was found to be accurate for low quark masses. We similarly find it to be a very accurate representation of the mass splitting up to a quark mass of roughly 1 GeV.

The dashed-dotted line corresponds to the mass splitting when the hyperfine interaction is set to zero and the strength of the confining term is increased so that the light constituent quark mass remains unchanged. It follows that even in the chiral limit roughly 80% of the $\pi - \rho$ mass splitting is due to the presence of the hyperfine interaction. This ratio is 92% and 99% for $K^* - K$ and $D^* - D$ respectively. We stress that the numerical value of this ratio depends on the details of the confining interaction. It is possible, for example, to obtain the correct $\pi - \rho$ mass splitting without a hyperfine interaction by adjusting the strength of the confining potential. The large m_Q behavior of the vector-pseudoscalar mass difference would however not be properly described, ie., it would decrease with m_Q more rapidly than $1/m_Q$. In this case it also found that the confining interaction alone leads to predictions for the quark condensate and pion decay constant which are too small. For example, typical predictions for the condensate are roughly -(95 MeV)^3 [16]; which should be compared with the phenomenological value of -(250 MeV)^3. In this model we obtain $\langle \bar{q}q \rangle(\Lambda_R) = -(200 \text{ MeV})^3$. The improvement is due to the hyperfine term, which has not been accounted for in previous calculations.

We now examine whether it is possible for the quark model to mimic the chiral behavior of the hyperfine splitting shown in Fig. 1. Our hypothesis is that the freedom to define quark masses in the quark model is sufficient to reproduce the splitting. We address this by plotting the hyperfine splitting as a function of the constituent quark mass derived from Eq. 3 (this is shown as a dashed line in Fig. 1). The curve reproduces the observed splitting for $\rho - \pi$, $K^* - K$, $D^* - D$ and $B^* - B$, shown as squares in the figure. We conclude that it is possible for a CQM to mimic the effects of chiral symmetry breaking. This is true because chiral symmetry breaking creates massive quasiparticles which may be used as effective degrees of freedom in model building. Furthermore, the hyperfine $1/m_Q$ behavior which is valid for heavy quarks continues to be valid for lighter constituent quarks.

We emphasize that the solid line in Fig. 1 is a prediction which may be compared with lattice data (triangles). The only adjustable parameters, the two renormalized potential strengths, are fixed by the pion and rho masses when using 5 MeV for the renormalized light quark mass.

We now consider hyperfine splittings in quarkonium systems as a function of the quark mass. In order to effectively mimic a Coulombic bound state for large m_R and a confined state for small m_R the renormalization scale is set to $\Lambda_R = \tilde{c}_R m_R + (m_R b)^{1/3} + b^{1/2}$. In contrast with the heavy-light open flavor system, the hyperfine splitting does not vanish with increasing heavy quark mass m_R but becomes proportional to $\tilde{c}_R^4 m_R$, as expected for a Coulombic system. We also considered the case where the renormalization scale corresponds to purely "confined" systems, *i.e.*, $\Lambda_R = (m_R b)^{1/3} + b^{1/2}$. This leads to $M_{hfs} \propto \Lambda_R^2/m_R \propto m_R^{-1/3}$, a behavior which is different from that expected in the constituent quark model (where $M_{hfs} \propto \Lambda_R^3/m_R^2 \propto m_R^{-1}$). The difference can be traced to the presence of negative energy solutions in the RPA equation. Elimination of the negative energy solutions leads to an effective hyperfine interaction which is proportional to $\tilde{c}_{R,h}^2 \Lambda_R^2/m_R$, while H_h alone leads to $\tilde{c}_{R,h}\Lambda_R^3/m_R^2$. Thus it is clear that correctly choosing the renormalization scale is crucial to establishing contact with quark model phenomenology.

We conclude with a discussion of the renormalization procedure. In the simple model studied here there are three renormalization scale-dependent quantities: $m(\Lambda)$, $\tilde{c}_c(\Lambda)$, and $\tilde{c}_h(\Lambda)$. We have used the pseudoscalar and vector meson masses to determine the (nonperturbative) renormalization scale dependence of these parameters. The third renormalization condition is obtained by fixing m_{con} (or alternatively $m(\Lambda_R)$ as done above for the light quarks) and by taking m_{con} to be Λ-independent. For large Λ/Λ_R Eq. (3) implies that the bare quark mass behaves as $m(\Lambda) \propto m_{con}^3 \ln(\Lambda/m_{con})/\Lambda^2$. Furthermore, at large cutoff the couplings saturate to constants. Thus asymptotic freedom is not recovered in our model. This is because the contact central interaction leads to matrix elements of order $\tilde{c}_c(\Lambda)$. In contrast, the full nonabelian Coulomb potential yields matrix elements of order $\tilde{c}_c(\Lambda)\ln(\Lambda)$. Asymptotic freedom is obtained because renormalization then requires that $\tilde{c}_c(\Lambda) \propto 1/\ln(\Lambda)$ for large Λ. Instead of choosing m_{con} to be a physical parameter one may use the condition $d[m(\Lambda)\langle\bar{q}q\rangle(\Lambda)]/d\Lambda = 0$. This leads to a running constituent mass $m_{con} = m_{con}(\Lambda)$ which vanishes for large Λ.

The use of the contact interaction rather than the full Coulomb interaction leads to several other differences in phenomenology. For example, the constituent quark mass becomes a function of the quasiparticle momentum when

the Coulomb potential is employed. Single particle energies are enhanced (even divergent) due to the confining nature of the nonabelian Coulomb potential. Thus, in a realistic calculation one cannot use the constituent quark energy (or constituent mass) to constrain parameters of the interaction. Instead one may, for example, use the condition on the quark condensate given above. Finally, the contact interaction leads to a single bound state per channel which certainly precludes analysis of the corresponding quark model excited states. While all of these points affect details of our computation, they do not change the main conclusion.

Spontaneous chiral symmetry breaking causes a pseudoscalar-vector meson mass splitting, but the hyperfine interaction is required to obtain the necessary $1/m_Q$ behavior for both heavy and light constituent quark masses. The interaction between the constituent quarks need not retain a memory of the underlying chiral symmetry and is described well by an effective hyperfine interaction. Thus Fig. 1 represents a direct validation of the main assumption of the naive quark model.

We have shown how quark model phenomenology may be derived from a simple model of QCD. Elimination of high momentum components from quark-transverse-gluon coupling leads to short range hyperfine interactions. By choosing a renormalization scale which matches the quark model we neglect contributions from long range interactions (due to the exchange of low energy gluons) to the hyperfine interaction – as is consistent with the quark model. Then, studying the quark mass dependence of the pseudoscalar-vector mass splitting we are able to show that interactions between constituent quarks derived from QCD indeed follow that of the naive quark model while respecting chiral symmetry. As a result we have shown how the heavy quark mass limit extrapolates to the chiral limit and have illustrated the interplay between hyperfine interactions and chiral dynamics. The potential of utilizing standard many-body techniques in applications to QCD in the Coulomb gauge has been discussed by the authors and many others [5,16,17]; however, neither the role of renormalization in constructing the Hamiltonian and building the constituent quark representation nor a QCD-based demonstration of the applicability of the effective quark model interactions to both heavy and light quarks has been addressed in this context before. Even though the analysis presented here has used a simplified central potential, the key results are independent of this choice. This is because the simplified model captures all features (*i.e.*, the presence of hyperfine interactions, correct momentum scales and dynamical chiral symmetry breaking) of QCD which are relevant. A study employing the full nonabelian potential will be presented elsewhere.

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